

# Upper bound on all products of R-parity violating couplings $\lambda'$ and $\lambda''$ from proton decay

Alexei Yu. Smirnov<sup>a,b</sup> and Francesco Vissani<sup>a,c</sup>

<sup>a</sup> *International Centre for Theoretical Physics, ICTP  
Via Costiera 11, I-34013 Trieste, Italy*

<sup>b</sup> *Institute for Nuclear Research,  
Russian Academy of Sciences,  
117312 Moscow, Russia*

<sup>c</sup> *Istituto Nazionale di Fisica Nucleare, INFN  
Sezione di Trieste*

## ABSTRACT

We prove that *any* product of R-parity violating couplings  $\lambda'$  (L-violating) and  $\lambda''$  (B-violating) can be strongly restricted by proton decay data. For any pair  $\lambda'$  and  $\lambda''$  the decay exists at least at one loop level. For squark masses below 1 TeV we find the conservative bounds  $|\lambda' \cdot \lambda''| < 10^{-9}$  in absence of squark flavor mixing, and  $|\lambda' \cdot \lambda''| < 10^{-11}$  when this mixing is taken into account. We study the dependence of the bounds on the flavor basis in which R-parity breaking couplings are determined.

1. The proton decay gives very strong bounds on the products of the R-parity violating couplings  $\lambda'$  (L-violating) and  $\lambda''$  (B-violating) involving *light* generations [1]. The decay takes place at tree level, and for squark masses about 1 TeV one gets [2]:

$$|\lambda' \cdot \lambda''| \lesssim 10^{-24}. \quad (1)$$

What are the bounds on the couplings involving heavy generations? Whether there are unrestricted couplings? These questions are important not only for phenomenology [3, 4, 5, 6, 7, 8, 9] but also for understanding the origin of the R-parity, as well as for the unification of interactions. Notice, for instance, that in stringy unified models the R-parity violation could be the only mechanism of the proton decay.

In the context of the Grand Unified Theories it was shown that the quark-lepton symmetry usually leads to separate strong bounds on  $\lambda'$  and  $\lambda''$  [10]. The situation can be different in absence of Grand Unification. It was claimed that certain products of  $\lambda'$  and  $\lambda''$  constants may be large without inducing proton decay at observable level [11].

In this Letter we prove that operators relevant to the proton decay are always present in the one loop effective lagrangian, and they imply strong bounds on any product of the couplings  $\lambda' \cdot \lambda''$ .

2. Let us introduce the L-violating coupling constant  $\lambda'_{ijk}$  and the B - violating coupling constant  $\lambda''_{mnp}$ , where  $i, j, k, m, n, p = 1, 2, 3$  are the generation indices, as the constants of the R-parity violating interactions:

$$\begin{aligned} & \lambda'_{ijk} D_i^{c\alpha} (\nu_j S_{kl}^d D_l^\alpha - E_j S_{kl}^u U_l^\alpha) \\ + & \lambda''_{mnp} \epsilon_{\alpha\beta\gamma} D_m^{c\alpha} D_n^{c\beta} U_p^{c\gamma} \end{aligned} \quad (2)$$

which are consistent with the Standard Model symmetry. Here,  $E_i, \nu_i, D_i^c, D_i, U_i^c, U_i$  are the superfields with charged leptons, neutrinos, down- and up-type-quarks, and  $\alpha, \beta, \gamma$  are color indices. Notice that in (2) the sum is over the flavor index  $l$  only. The couplings are written in terms of superfields whose fermionic components coincide with the mass eigenstates (in all orders of perturbation theory). The unitary matrices  $S^d$  and  $S^u$  connect the fermionic states of the original basis in which the couplings are defined with the mass states. The product

$$S^{u\dagger} S^d = V \quad (3)$$

gives the Cabibbo-Kobayashi-Maskawa matrix  $V$ .

Let us stress that the R-parity violating couplings depend on separate rotations of the upper,  $S^u$ , and down,  $S^d$ , components of the quark doublet. Correspondingly, the bounds on the couplings will depend on the basis in which they are defined. One can also introduce flavor rotations  $S$  for the other quark and lepton superfields.

The scheme of the proof that proton decay is induced by any pair of the couplings  $\lambda', \lambda''$  is the following. Using an arbitrary pair  $\lambda'$  and  $\lambda''$ , one can construct 4-field effective operators with (B - L) or (B + L) violation. If only light quarks and leptons are involved, these operators leads to the proton decay already at tree level. If the operator contains heavy quarks the proton decay will be forbidden kinematically. However, additional interactions with charged Higgs bosons (or Higgsinos,  $W$ -bosons or Wino), which violate flavor and thus can transform heavy quarks into light ones, lead to operators inducing proton decay. Furthermore we will show that for all pairs of couplings such an operator appears at the one loop level.

The coupling of the physical charged Higgs boson  $h^+$  with the quarks and the squarks can be written as:

$$V_{kl} h^+ \left\{ \frac{m_{d_l}}{v} \left[ \tan \beta \cdot u_k d_l^c - (\mu + A_{d_l} \tan \beta) \cdot \tilde{u}_k \tilde{d}_l^c \right]^* + \frac{m_{u_k}}{v} \left[ \cot \beta \cdot u_k^c d_l - (\mu + A_{u_k} \cot \beta) \cdot \tilde{u}_k^c \tilde{d}_l \right] \right\} + \text{h.c.} \quad (4)$$

The tilded fields correspond to the scalar quarks,  $A_{u,d}$  are the soft breaking parameters,  $\mu$  is the supersymmetric mass,  $\tan \beta$  is the ratio of vacuum expectation values of the two Higgs fields and  $v = 174$  GeV. Notice that the up-quark masses,  $m_{u_k}$ , appear in the interactions of  $u_k^c$  and  $\tilde{u}_k^c$  fields, whereas the interactions of the  $d_l^c$  and  $\tilde{d}_l^c$  are proportional to the down-quark masses  $m_{d_l}$ .

The squarks  $\tilde{d}_n$  and  $\tilde{d}_n^c$  mix by the mass:

$$m_{\tilde{d}_n}^2(\text{LR}) = m_{d_n}(A_{d_n} + \mu \tan \beta), \quad (5)$$

and similarly will do the  $u$ -type-squarks.

3. In general interfamily connections enhance the proton decay. Therefore to get the most conservative bound one should use the basis with maximally suppressed connection. For this purpose we choose the basis in which

$$S^d = I, \quad S^u = V^\dagger. \quad (6)$$

In the proof we will use the L-violating coupling with neutrino only (the first term in (2)) which does not contain any mixing at all. We will discuss the dependence of the result on the basis in Sect. 6.

The generation structure of  $\lambda'$  and  $\lambda''$  plays a crucial role. Let us recall in this connection that, according to Eq. (2), in the  $\lambda'$ -coupling we use the index  $i$  to denote the generation of  $D^c$ , the index  $j$  for the lepton doublet and  $k$  for the quark doublet. In the  $\lambda''$ -coupling the indices  $m$  and  $n$  are prescribed to the  $D^c$  superfields, whereas the index  $p$  is for the  $U^c$  superfield.

All possible pairs of couplings  $\lambda'_{ijk}$ ,  $\lambda''_{mnp}$  can be divided into two classes:

- (i) Pairs with “matching” of the  $D$ -fields, when the generation index of the  $D$  or  $D^c$  field from the lepton-violating vertex coincides with the generation index of the  $D^c$  field from the baryon-violating vertex. That is  $i = m$  or  $i = n$  (“ $D^c D^c$ -matching”), or  $k = m$  or  $k = n$  (“ $D^c D$ -matching”).
- (ii) Pairs without “matching” of the  $D$ -fields. Taking into account the antisymmetry of  $\lambda''_{mnp}$  in  $m$  and  $n$  one finds that “no-matching” case is realized only if  $i = k$  and  $i, m, n$  are all different, *i.e.*  $D^c$ -fields of all three generations should be present. Let us construct and estimate the diagrams for these two cases.

- (i) Pairs with “matching”.

Suppose first that there is the “ $D^c D^c$ -matching” and take  $i = m$  for definiteness. We can connect the B- and the L-violating vertices by one  $\tilde{d}^c$ -propagator. In the obtained tree level diagram (Fig. 1a) at most one external line is  $b$  or  $b^c$ . (If both of them are  $b$ , then one has situation with “ $D^c D$ -matching” which will be considered later). The case  $d_n^c \equiv b^c$  corresponds to  $d_k = d$  or  $s$ ; the former quark should be connected with  $u_p^c$  by charged Higgs, whereas the latter can be an external line of an operator which gives proton decay. In such a way one gets the vertex type diagram shown Fig. 1c. Emitting  $h^+$ ,  $b^c$  transforms into an  $u$ -quark, and, absorbing  $h^+$ ,  $u_p^c$  transforms into a  $d$  (or a  $s$ ) quark. The resulting effective operators,  $udd\nu$  or  $udsv$ , contain all light fields. Using the coupling constants of the Higgs (4) we find the suppression factor,  $\xi$ , of the one loop diagram with respect to the tree level diagram  $\xi \equiv (loop)/(tree)$ :

$$\xi \approx \frac{1}{(4\pi)^2} \times \frac{m_{u_p}}{v} \frac{m_b}{v} \times V_{13} V_{p1} . \quad (7)$$

If  $d_k \equiv b$ , then one should connect by Higgs exchange  $u_p^c$  with  $d_k$ , so that  $d_k \rightarrow u$  and  $u_p^c \rightarrow d$  or  $s$ . The resulting box diagram, of the type shown in Fig. 1e, has a suppression

factor similar to that in (7) with substitution  $m_b V_{13} \rightarrow m_u V_{1k}$ .

Suppose that in the tree-level diagram both  $d$ -lines are not  $b$ ; the decay is allowed unless  $u_p^c$  are  $c^c$ - or  $t^c$ -quarks, or if both  $d$ -lines are  $s$ -quarks. In both cases, the Higgs exchange between a  $d^c$  line and  $u_p^c$  leads to the proton decay at one loop level, and the suppression factors are similar to that in Eq. (7).

In the case of “ $D^c D$ -matching” among the external lines in the diagram (Fig. 1b) at most one is  $b$  and at least one coincides with  $d$  or  $s$ . Similarly to the previous case one should connect this heaviest  $d^c$ -quark line with  $u_p^c$ , thus arriving at vertex (Fig. 1d) or box (Fig. 1f) diagrams. The suppression factor equals

$$\xi \approx \frac{1}{(4\pi)^2} \times \frac{m_{d_n}}{v} \frac{m_{u_p}}{v} \frac{m_b}{\tilde{m}} \times V_{1n} V_{1p} , \quad (8)$$

where  $\tilde{m}$  is the typical squark mass and the factor  $m_b/\tilde{m}$  follows from the mixing of the left and right squarks (5). Notice that the diagrams shown lead to (B+L) conserving operators  $udd\bar{\nu}$ ,  $uds\bar{\nu}$ .

(ii) “No-matching” case.

Recall that in “no-matching” case the three  $D^c$  should all be of different flavors and therefore one of them is  $B^c$  and two others are  $D^c$  and  $S^c$ . In this case one can always construct the diagram of the type shown in Fig. 2. The  $\tilde{u}_p^c$ -squark from the baryon-violating vertex can emit an Higgs boson and a  $\tilde{d}_k$ -squark, and the latter can be adsorbed by the lepton-violating vertex. The higgs field in turn is absorbed by the  $d_i^c \equiv b^c$  line, so that  $b^c \rightarrow u$  (Fig. 2a, 2b). As a result one gets the operator  $udsv$ . Integration over the loop, and the Lorentz structure of the vertices, pick up the momentum of an external quark  $p_q \sim m_N/3$ , where  $m_N$  is the nucleon mass. The suppression factor can be estimated as

$$\xi \approx \frac{1}{(4\pi)^2} \times \frac{m_b}{v} \frac{m_{u_p}}{v} \frac{p_q}{\tilde{m}} \times V_{31} V_{pk} \tan \beta. \quad (9)$$

For a given couplings one can construct another version of the diagram: the field  $b^c$  changes chirality,  $b^c \rightarrow b$ , and  $b$  transforms into  $u^c$  by absorbing an Higgs field (Fig. 2c, 2d). In this case

$$\xi \approx \frac{1}{(4\pi)^2} \times \frac{m_u}{v} \frac{m_{u_p}}{v} \frac{m_b}{\tilde{m}} \times V_{31} V_{pk} \tan \beta. \quad (10)$$

(Notice that now  $p_q/\tilde{m}$  is substituted by  $m_u/v$ ). The flip of chirality may take place in  $\tilde{d}-$  (Fig. 2e) or  $\tilde{u}-$  (Fig. 2f) propagators connecting the B- and the L-vertices. The suppression factors are of the same type as in (10).

Thus we have shown that for any pair  $\lambda', \lambda''$  there is (at least one) one loop diagram which leads to the proton decay. Usually for each pair one can find several diagrams, and moreover in some cases, as we will see in Sect. 4, new diagrams give even bigger contributions.

Diagrams similar to those in Figs. 1, 2 arise due to the  $W$ -boson, the would-be Goldstone boson and the charginos (mixed states of charged Higgsinos and Wino) exchanges, instead of the Higgs exchange. Effective tree level or one loop operators leading to the emission of light charged leptons are possible for certain pairs of couplings in (2). Additional sources of flavor violation could be related to the interactions of gluinos and neutralinos.

For each channel of the proton decay one can find different contributions. Since in general they have different Lorenz structure and depend on different unrelated parameters (masses *etc.*), we suggest that there is no accidental strong cancellation of different contributions.

4. Let us find now a conservative bound on the product of the R-parity violating couplings and identify the pairs of coupling which are less restricted. For this we evaluate the suppression factors  $\xi$  corresponding to each constructed diagrams. We assume squark masses around 1 TeV, and take quark running masses at the squark scale. Since for large values of  $\tan \beta$  the suppression is weaker, we will use  $\tan \beta = 2$ , compatible with the top Yukawa coupling being perturbative up to the Planck scale.

Considering the diagrams Fig. 1,2 only, we found that the smallest factor  $\xi$  is for the pair of couplings

$$\lambda'_{3j3} \lambda''_{121} \quad (11)$$

(the L-violating and the B-violating vertices contains  $B^c \nu_j B$  and  $D^c S^c U$  correspondingly). There is no “ $D$ -matching”. Dominating contribution comes from “no-matching” diagram of Fig. 2a. The suppression factor is given in (9) with  $V_{pk} = V_{13}$ :  $\xi \approx 10^{-17}$ . The factor results in rather weak bound on the couplings  $|\lambda'_{3j3} \cdot \lambda''_{121}| < 10^{-7}$ . However for these couplings another type of diagrams exists, the one with neutrino-Zino mixing, which leads to a much stronger bound.

The neutrino and the Zino are mixed by one loop diagrams generated, *e.g.*, by the R-parity violating coupling  $\lambda'_{3j3} b^c \nu \tilde{b}$ , the gauge interaction  $g \tilde{b}^* b \tilde{Z}$  and by the mass term

$m_b \, bb^c$ . The Zino couples with one of the squarks emitted from the B-violating vertex thus leading to the proton decay. Notice that for the existence of such a diagram it is important that all the quarks in the B-violating vertex are light. The suppression factor for the diagram with  $\nu - \tilde{Z}$  mixing can be estimated as

$$\xi \approx \frac{g^2}{(4\pi)^2} \times \frac{m_b}{m_{\tilde{Z}}} \approx 7 \cdot 10^{-6}. \quad (12)$$

It gives the bound on the product of couplings of the order  $10^{-18}$ .

Next weakest bound is for the product

$$\lambda'_{2j2} \lambda''_{131}. \quad (13)$$

For these couplings there is no proton decay diagram with  $\nu - \tilde{Z}$  mixing. The dominating contribution comes from the “no-matching” diagram of Fig. 2b, and the suppression factor is:

$$\xi \approx 10^{-15}. \quad (14)$$

Other products,

$$\lambda'_{1j2} \lambda''_{231}, \quad \lambda'_{3j2} \lambda''_{121}, \quad \lambda'_{1j1} \lambda''_{231}, \quad \lambda'_{2j1} \lambda''_{131}, \quad \lambda'_{3j1} \lambda''_{121}, \quad (15)$$

have a few times larger  $\xi$  factors. All but the third pair in (15) correspond to the case of “ $DD^c$ -matching.” but since the “matching” is suppressed by the LR mixing the less suppressed diagrams are again those of Fig. 2.

We conclude, using (14), that the conservative bound on *any* product of  $\lambda'$ - and  $\lambda''$ -type coupling is:

$$|\lambda' \cdot \lambda''| \lesssim 10^{-9}. \quad (16)$$

5. Let us consider the effect of flavor-changing squark mixing. The mixing is induced by the Yukawa interactions in the superpotential and by the corresponding soft symmetry breaking terms. The mixing of the right handed components  $\tilde{b}^c - \tilde{s}^c$  and  $\tilde{b}^c - \tilde{d}^c$  proceeds via one loop diagrams formed by  $u_p \tilde{h}$  or  $\tilde{u}_p h$ , or by loops of  $h$  and  $\tilde{u}_p$  connected to  $\tilde{b}^c - \tilde{s}^c$  by four boson coupling. Summation over  $p$  leads to GIM cancellation which renders finite the contribution, and the suppression factor can be estimated as

$$\xi_{bs} \approx \frac{1}{(4\pi)^2} \times \frac{m_b m_s}{v^2} \frac{m_t^2}{\tilde{m}^2} \times V_{23} V_{33} \approx 3 \cdot 10^{-8}. \quad (17)$$

For  $\tilde{b}^c - \tilde{d}^c$  mixing  $m_s V_{23}$  should be substituted by  $m_d V_{13}$ .

In the case of the left-right type mixing:  $\tilde{b}^c - \tilde{s}$  and  $\tilde{b}^c - \tilde{d}$  the  $u_p$  quark in the loop should change chirality:  $\tilde{b}^c \rightarrow t\tilde{h} \rightarrow t^c\tilde{h} \rightarrow \tilde{s}$ . There is no GIM cancellation and the diagrams are logarithmically divergent. A typical suppression factor equals

$$\xi_{bs} \approx \frac{1}{(4\pi)^2} \times \frac{m_b m_t}{v^2} \frac{m_t m_{\tilde{h}}}{\tilde{m}^2} \times V_{23} V_{33} \times \ln \frac{M_P}{\tilde{m}} \approx 3 \cdot 10^{-6}, \quad (18)$$

where the renormalization point has been chosen at the Planck scale  $M_P$ . For  $\tilde{b}^c - \tilde{d}$  mixing the suppression factor can be obtained from (18) by the substitution:  $V_{23} \rightarrow V_{13}$ .

The suppression factors (17, 18) are consistent with the bounds on the flavor-changing neutral currents, and there is no reason to disregard flavor-changing squark mixing.

The flavor-changing mixing is important for restrictions on the couplings with large number of light generation indices. New tree level diagrams appear with squark mixing in the propagator. In fact, for all pairs of the couplings (13,15) one can construct such a kind of diagram. This leads to suppression factors of the order  $10^{-8} - 10^{-6}$  instead of  $10^{-15}$  found previously in absence of mixing. Correspondingly, one gets very strong bounds  $|\lambda' \cdot \lambda''| < 10^{-16}$ .

What are the largest allowed products of couplings in this case? Tree level diagrams with squark mixing for proton decay are absent for pairs of couplings with (1) two or more unmatching heavy fields (which can not be produced for kinematical reason) or (2) with more than two  $s$ -fields. For example, the first criteria is satisfied for couplings with  $c^c$ - or  $t^c$ -quark-fields instead of  $u^c$ . The smallest suppression factors are found for the pairs

$$\lambda'_{1j1} \lambda''_{232}, \quad \lambda'_{2j1} \lambda''_{132}, \quad \lambda'_{3j1} \lambda''_{122}, \quad \lambda'_{3j3} \lambda''_{131}, \quad \lambda'_{3j3} \lambda''_{231}, \quad (19)$$

where  $j = 1, 2, 3$ . The main contributions to the proton decay come from the following diagrams: pairs 1,2—"no-matching" type vertex diagram 2b; pair 3—"B<sup>c</sup>B-matching" vertex, diagram 1d; pairs 4,5—"B<sup>c</sup>B-matching" box diagram of Fig. 1e. The suppression factor  $\xi \approx 10^{-13}$  leads to the bound on the products of couplings (19):

$$|\lambda' \cdot \lambda''| \lesssim 10^{-11}. \quad (20)$$

One remark is in order. The tree level diagram for the above couplings contains only one heavy quark  $b$ . In general due to the wave function renormalization this quark can be transformed into a  $d$  or a  $s$  leading to very fast proton decay. However such a renormalization is absent for our definition of the couplings: Recall that they are introduced in



the basis where the fermionic components of the supermultiplets coincide with the mass states in any order of perturbation theory. This is equivalent to the redefinition of the couplings in each order of the perturbation theory.

6. Let us consider the dependence of the bounds on basis in which the couplings are defined. As we marked before, the R-parity breaking couplings depend on the rotation matrices  $S^u$  and  $S^d$  separately, as well as on the possible rotation of the  $D^c, U^c, L$  and  $E^c$  superfields. For example, in the basis where  $S^d = V$  and  $S^u = I$ , there is “matching” for all pairs of couplings.

In general, the bounds in any basis can be obtained from the bounds in the original basis by an appropriate rotation. Let us show that the conservative bound (20) is in fact basis independent.

In a basis characterized by  $S^d \neq I$  the coupling constants  $\lambda'_{ijk}$  are related to the constants in the original basis,  $\lambda_{ijk}^{0'}$ , as

$$\lambda'_{ijl} S_{lk}^d = \lambda_{ijk}^{0'}. \quad (21)$$

Since  $S_{lk}^d$  is an unitary transformation we get from the bound  $|\lambda_{ijl}^{0'}| < \mathcal{B}_0$ :

$$\sum_k |\lambda'_{ijk}|^2 = \sum_l |\lambda_{ijk}^{0'}|^2 < 3 \mathcal{B}_0^2, \quad (22)$$

where  $\mathcal{B}_0$  is the conservative bound on all modes in the original basis. Similarly one can estimate the effect of the rotation of the other superfields.

7. Final remarks. We have shown that for any pair of R-parity violating couplings it is possible to construct one loop diagrams which lead to the proton decay.

If one disregards the flavor-changing squark mixing then the most conservative bound on the product of couplings is  $|\lambda' \cdot \lambda''| \lesssim 10^{-9}$  for squark masses below the TeV scale. If one takes into account flavor-changing squark mixing then new diagrams appear and the conservative bound becomes stronger:  $|\lambda' \cdot \lambda''| \lesssim 10^{-11}$ .

The pairs of coupling which have the weakest suppression are given in (13), (15) and (19). It is interesting to observe that the couplings involved are not those with maximal number of the heavy generation indices. In fact these couplings contain typically several light indices. Therefore it may be difficult to explain the dominance of these couplings by some horizontal symmetry, especially in models which reproduce the hierarchy of the

fermion masses. Let us admit however that some mechanism exists which picks up a pair of constants which are weakly suppressed. In this case the renormalization of these  $\lambda$  due to Yukawa interactions will induce R-parity violating couplings which are not weakly bounded. To solve this problem one should suggest that this mechanism operates at the scale  $\mu \sim \tilde{m} \sim 1$  TeV. This mechanism will not operate (or will imply further fine-tuning) if  $\mu$  is of the order of the Grand Unification scale  $M_{GU}$  or of the Planck scale.

The bounds (16) and (20) are considerably weaker than the bound (1), and weaker than the typical bound which can be obtained in the context of Grand Unified theories. The bound (16) is marginally compatible with both lepton *and* baryon violating effects which could be observable at accelerators,  $\lambda_{\text{obs}} \gtrsim 2 \cdot 10^{-5} \sqrt{\gamma} (\tilde{m}/1 \text{ TeV})^2 (150 \text{ GeV}/m_\chi)^{5/2}$ , where  $\gamma$  is the Lorentz boost factor [9]. For example, a neutralino  $\chi$  of approximately 150 GeV (or heavier) mass decays within the detector by couplings  $\lambda' \sim \lambda'' \sim 2 \cdot 10^{-5}$  which satisfy the proton decay bound (16). However, according to previous discussion, this case would require various fine-tunings of the model.

The bounds discussed in the present paper will be strengthened at least by 1 order of magnitude if the SuperKamiokande experiment will not detect a signal of proton decay.

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## Figure Captions

Fig. 1: Feynman diagrams inducing proton decay in the (anti)neutrino-meson(s) channels for the “matching” case, but not for the “no-matching” case. The blob indicates a fermion mass insertion.

Fig. 2: Feynman diagrams inducing proton decay in the (anti)neutrino-meson(s) channels for the “no-matching” case. Compare with previous figure.

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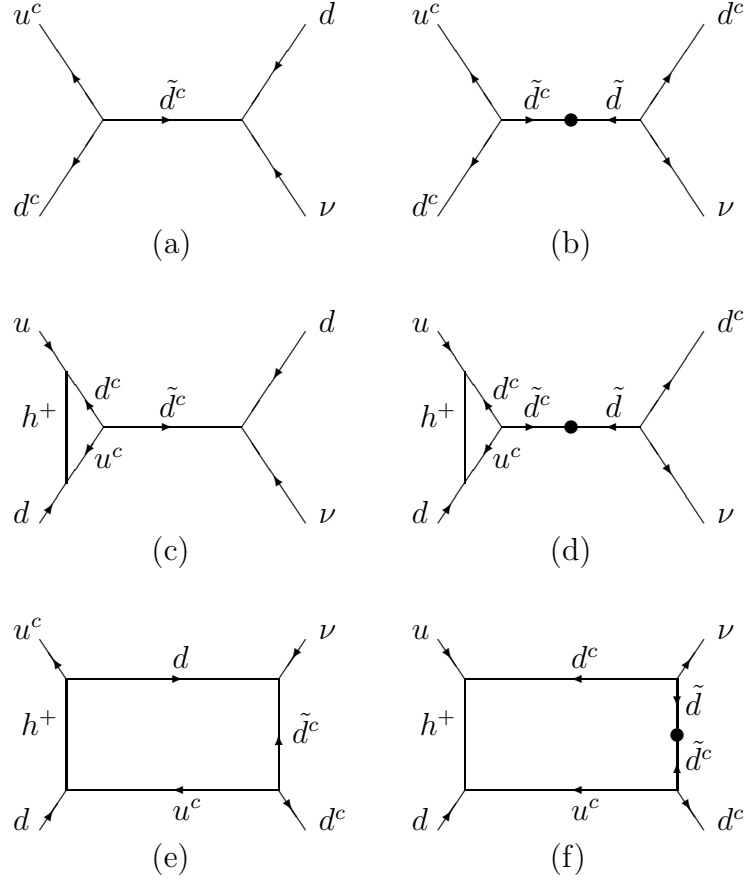


Figure 1: Feynman diagrams inducing proton decay in the (anti)neutrino-meson(s) channels for the “matching” case, but not for the “no-matching” case. The blob indicates a fermion mass insertion.

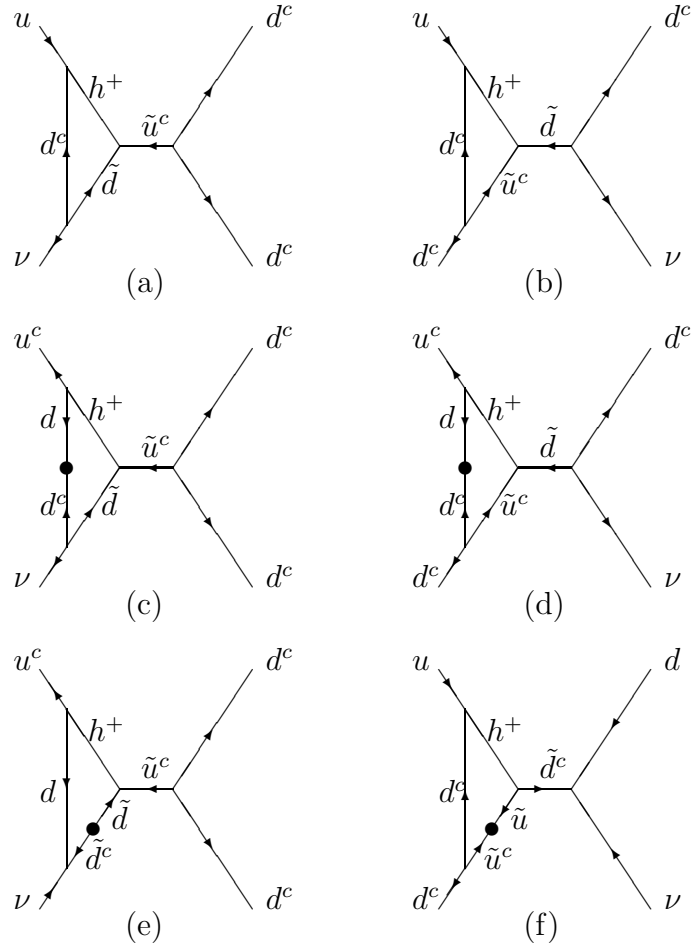


Figure 2: Feynman diagrams inducing proton decay in the (anti)neutrino-meson(s) channels for the “no-matching” case. Compare with previous figure.